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Chirality, D-branes, and Gauge/String Unification

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We point out that models where gauge and charged matter fields live on the world volume of D3 branes can provide a weak string coupling description of the gauge coupling unification in the supersymmetric standard model at a scale of order 10^{16} GeV. In the D3 brane picture, the string and unification scales coincide, while the size of the extra dimensions is somewhat larger than the string length. We then construct a quasirealistic "three generation" model with D3 branes at orbifold fixed points. We point out that placing D3 branes at different orbifold fixed points provides a geometric description of the disconnected branches of moduli space and discuss the appropriate consistency conditions. These disconnected branches allow for patterns of gauge symmetry breaking inaccessible on the connected part of moduli space and can give rise to "hidden" and "visible" sectors. Finally, we analyze supersymmetry breaking on the brane world volume. (Based on a talk presented at the XXIIIrd Rencontres de Moriond "*Electroweak Interactions and Unified Theories*," Les Arcs, Savoie, France, March 14-21, 1998, and work with J. Lykken and S.P. Trivedi, to appear).

1. Introduction and summary.

In the past several years, we have witnessed a tremendous progress of our understanding of string theory. The five string theories, previously thought to be different, have been shown to be related by a complicated web of dualities [1]. A crucial role in this understanding is played by the extended nonperturbative objects in string theory, known as D-branes [2]. Dp-branes are p spatial dimensional hypersurfaces embedded in ten dimensional space time, whose world volumes support Yang-Mills theories with varying amount of supersymmetry. Furthermore, constructions using various branes give a unified description of many different gauge theories in various dimensions [3] and provide some fascinating connections between string and gauge theory dynamics, which are still being unraveled.

Superstring theory gives rise to gravity, gauge, and matter (fermion) fields, and is thus is a natural candidate for a theory unifying the known forces. Most model-building efforts involving strings have, so far, been in the framework of the heterotic string theory. Since all string theories are related by dualities, one may ask whether the weakly coupled description of the low-energy physics could be in terms of one of the other, dual, string theories. It is therefore interesting to investigate the model-building possibilities in the other string theories.

Another motivation (admittedly, relying on certain assumptions) to look into different possibilities is provided by the fact that the weakly coupled heterotic string has difficulties accommodating the "observed" unification of gauge couplings [4] in the (supersymmetric) standard model at a scale 10^{16} GeV (for a review see [6]). Witten proposed [5] to consider instead the strong coupling dual of the $E_8 \times E_8$ string—the eleven-dimensional M-theory compactified on a segment (S^1/Z_2) times a six-dimensional compact manifold (this "M-theory phenomenology" has been the subject of some interest lately [7]). Ref. [5] also mentioned the possibility that Type I theories could give a weak coupling description of the gauge unification as well. Here we will elaborate somewhat on this issue and will point out some, hopefully, generic features of Type I "realistic" models.

We will concentrate on models, where the gauge and matter fields of the standard model live on the world volume of three spatial dimensional extended objects (D3-branes) embedded in ten dimensional space time. Since the closed strings (gravity) can also propagate in the bulk, in order to avoid violations of Newton's law at large

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distances the six dimensions transverse to the D3 branes have to be compactified (a discussion on experimental limits on the sizes of the extra dimensions is given in ref. [17]). For recent related work on the subject, see refs. [18], [19].

This paper is structured as follows. In Section 2, we begin with a discussion of gauge unification in models with D3 branes. We will see that the "observed" unification of couplings in the supersymmetric extensions of the standard model can be accommodated in models with D3 branes. The GUT and string scale are identified in these models, while the size of the extra dimensions is somewhat larger than the inverse GUT scale. Then, in Section 3, we consider a simple compactification of Type I theory with D3 branes, considered in [10] and more recently in [11]. We point out that, in addition to the connected part of moduli space, considered in [10], this model exhibits disconnected branches of moduli space. These branches are best visualized in terms of placing a set of D3 branes at orientifold planes away from the origin and are like the the disconnected vacua discussed in [9], [13]. We also briefly discuss the consistency conditions [12] that these branches have to obey. Along one of these branches, we obtain an $SU(5)$ model with three generations of $\bar{\mathbf{5}}$ and $\mathbf{10}$ matter fields. The model is not realistic—there are no Higgs fields to break the either the $SU(5)$ or the standard model gauge groups, and there are baryon and lepton number violating Yukawa couplings. However, we view the ease with which this "quasirealistic" matter content can be obtained in Type I compactifications as encouraging further study (for recent work, see [19]). Finally, in Section 4, we analyze in some detail the infrared dynamics of the $SU(5)$ three generation model. We show that for fixed finite values of the dilaton and the orbifold blow-up parameter the model dynamically breaks supersymmetry.

2. Gauge-string unification on D3 branes.

We will show here that theories with gauge and matter fields living on the world volume of D3 branes can, at weak string coupling, accommodate the "observed" unification of couplings of the (supersymmetric) standard model at a scale

of order 10^{16} GeV. We will see that in models with D3 branes, the string and "GUT" scales are naturally identified.[2] We will see that in the D3 brane picture, the extra dimensions are larger than the string scale and open up (in energy units) somewhat below the string scale. They do not, however, affect the running of the gauge couplings, since the gauge and charged matter fields are constrained to live on the world volume of the D3 branes. At the string scale, new charged states appear in the spectrum—the massive states of the open strings. The winding open string states are somewhat heavier than the string scale in this picture.

To find the relation between the string scale, string coupling, the size of the compact dimensions, and the Planck scale, consider the closed string effective action in 10 dimensions (see, e.g. [2]):

$$S_{grav,10d} = \frac{1}{2} \frac{1}{64\pi^7 \alpha'^4 g_s^2} \int d^{10}x \mathcal{R} + \dots, \qquad (1)$$

where $2\pi\alpha'$ is the string tension, and g_s—the string coupling. Upon dimensional reduction to 4 dimensions on a six-torus of volume $(2\pi R)^6$ this becomes:

$$
\begin{aligned}
S_{grav,4d} &= \frac{1}{2} \frac{(2\pi R)^6}{64\pi^7 \alpha'^4 g_s^2} \int d^4x \mathcal{R} \\
&= \frac{1}{2} \frac{M_{Pl}^2}{8\pi} \int d^4x \mathcal{R} + \dots, \qquad (2)
\end{aligned}
$$

where $\sqrt{M_{Pl}^2/8\pi} = 2.4 \times 10^{18}$ GeV is the reduced Planck mass. From eq. (2), we obtain the relation between the string scale, string coupling, and the radius of the compact dimensions:

$$\frac{R^3}{\sqrt{\pi}\alpha'^2 g_s} = 2.4 \times 10^{18} \text{ GeV} . \qquad (3)$$

The gauge and charged matter fields in models with Dp-branes are, on the other hand, constrained to live on the brane world volume. Hence, to find the relation between the string coupling and scale, and the gauge coupling, we need

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[2]Hereafter, by the "grand unification" scale, M_{GUT}, we mean the scale where the gauge couplings become equal to one another; one does not have to assume that there is a grand-unified gauge group.

to consider the world-volume theory, described by the Born-Infeld action [2]:

$$S_{BI} =$$

$$-\frac{1}{(2\pi)^p \alpha'^{\frac{p+1}{2}} g_s} \int d^{p+1}\sigma \sqrt{\det\left(G_{ab} + 2\pi\alpha'\ F_{ab}\right)}$$

$$= -\frac{(2\pi\alpha')^2}{(2\pi)^p \alpha'^{\frac{p+1}{2}} g_s} \frac{1}{4} \int d^{p+1}\sigma F_{ab}F^{ab} + \dots$$

$$= -\frac{1}{4g_{YM,p}^2} \int d^{p+1}\sigma F_{ab}F^{ab} + \dots . \qquad (4)$$

Here G_{ab} is the induced metric on the brane world volume, F_{ab} is the gauge field strength, and we have kept only the terms describing the Yang-Mills fields. From eq. (4), we conclude that, on a Dp brane, the Yang-Mills gauge coupling is related to the string coupling by

$$g_{YM,p}^2 = g_s (2\pi)^{p-2} (\alpha')^{\frac{p-3}{2}} . \qquad (5)$$

This relation is valid at the string scale, $(\alpha')^{-1/2}$, which is the cutoff scale of the effective world-volume theory.[3]

From now on we consider D3 branes. The gauge coupling on the D3 brane at the string scale, eq. (5), is:

$$\alpha_{YM} = \frac{g_{YM}^2}{4\pi} = \frac{g_s}{2} . \qquad (6)$$

We take the values of the unification scale $M_{GUT} \sim 10^{16}$ GeV and the gauge coupling, $\alpha_{GUT} = 1/25$, for the supersymmetric standard model [4]. From (6) we thus obtain a value for the string coupling $g_s \sim .08$, which is well within the weak coupling regime. On the other hand, taking $\alpha' \sim M_{GUT}^{-2}$ from eq. (3), we obtain $RM_{GUT} \simeq 3$. We see that identifying the string scale with the GUT scale in models with D3 branes is natural—the size of the extra dimensions is larger than the string (and GUT) scale (which, as can be easily seen, is not the case in the T-dual 9-brane picture). New states appear in the effective theory only at the string scale: the

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[3]We note that for nonabelian fields, the relation (5) can be modified by a normalization constant; it can be obtained by T-duality from the Type I 10-dimensional effective action, as in [1] (see also [8]).

massive open and closed string excitations have mass of order $(\alpha')^{-1/2}$. The winding open string modes—the lightest excitations of open strings beginning and ending on the D3 branes, but winding around the compact directions—on the other hand, are somewhat heavier than the string scale, $M_{winding} \sim RM_{GUT}^2 \sim 3M_{GUT}$.

3. A "three generation" chiral model on D3 branes.

In this section, we present a simple "three generation" model with D3 branes on orbifold singularities (the orbifold projection will be useful, among other things, to obtain a chiral gauge theory on the brane world-volume). The model is, admittedly, not a realistic model, but it will serve the purpose of making several generic points quite explicit. In addition to providing a weak string coupling description of the "observed" unification of couplings, the D3 brane picture has other potential benefits.

We will see that the D3 brane picture allows for a geometric description of the disconnected branches of moduli space of their world-volume theories. These disconnected branches of moduli space correspond to placing D3 branes at orbifold fixed points other than the origin (such fixed points only exist in compact orbifolds) and are like the ones discussed in [9], [13]. The disconnected branches of moduli space can have multiple uses. We will see below that they exhibit patterns of gauge symmetry breaking that are not possible on the connected part of moduli space. In addition, branes placed at different orbifold fixed points can serve as "visible" and "hidden" sectors; the latter can be responsible for supersymmetry breaking. The lightest excitations of strings stretching between branes at different fixed points transform as fundamentals under both the "hidden" and "visible" gauge groups; they may be instrumental in the communication of supersymmetry breaking.

As a simple example to illustrate the above points, we consider the $T^6/\mathbb{Z}_3$ orientifold in some detail. We take the D3 branes to stretch along $X^{1,2,3}$ and introduce complex coordinates, $z_1 = X^4 + iX^5, z_2 = X^6 + iX^7, z_3 =$

$X^8 + iX^9$, in the transverse six-dimensional space. The orientifold group is given by $G = \{1, \alpha, \alpha^2, \Omega R(-1)^{F_L}, \Omega R(-1)^{F_L}\alpha, \Omega R(-1)^{F_L}\alpha^2\}$, where $\alpha = e^{2\pi i/3}$, and the action on the transverse coordinates is

$$(z_1,\ z_2,\ z_3)\ \rightarrow (\alpha z_1,\ \alpha z_2,\ \alpha z_3)\ . \tag{7}$$

Ω denotes world-sheet orientation reversal, and R is a reflection $z_i \rightarrow -z_i, i = 1, 2, 3$ in the space transverse to the D3 branes. F_L is an operator that flips the sign of the left-moving Ramond states. In addition, the orientifold group acts on the Chan-Paton indices of the open string states stretching between different D3 branes. The action of the group elements, $\Omega R(-1)^{F_L}$ and α, on the Chan-Paton factors λ can be represented by the matrices:

$$\lambda \rightarrow \gamma_{\Omega R(-1)^{F_L}}\ \lambda^T\ \gamma_{\Omega R(-1)^{F_L}}^{-1}\ , \tag{8}$$

and

$$\lambda \rightarrow \gamma_\alpha\ \lambda\ \gamma_\alpha^{-1}. \tag{9}$$

The matrices γ_α and $\gamma_{\Omega R(-1)^{F_L}}$ must furnish a representation of the orientifold group. The matrices $\gamma_{\Omega R(-1)^{F_L}}$, representing the action of the $\mathbb{Z}_2$ part of the orientifold group should obey

$$\gamma_{\Omega R(-1)^{F_L}}\ = \left(\gamma_{\Omega R(-1)^{F_L}}\right)^T\ . \tag{10}$$

In the absence of the $\mathbb{Z}_3$ orbifold projection, the $\Omega R(-1)^{F_L}$ projection would lead to an SO gauge group on the D3 brane world volume.

The untwisted Ramond-Ramond 4-form charge conservation conditions require the presence of 32 D3 branes to cancel the orientifold charge. On the other hand, the charge cancellation conditions for the twisted RR fields result in the following requiremant on the matrices γ_α:

$$\text{Tr}\ \gamma_\alpha\ =\ -4\ . \tag{11}$$

This condition should be imposed on D3 branes placed at fixed points of $\mathbb{Z}_3$ (i.e. at the origin and the two other $\mathbb{Z}_3$ fixed points in the interior of T^2).

In Fig. 1 we show one of the three two-tori of this orbifold ($T^6 = (T^2)^3$). On the compact T^2 (as opposed to a noncompact $\mathbb{C}$), there is

more than one fixed point of the orientifold $\mathbb{Z}_6$ group. The origin, shown as a triangle on Fig. 1, is the only $\mathbb{Z}_6 = \mathbb{Z}_2 \times \mathbb{Z}_3$ fixed point. In addition, there are three $\mathbb{Z}_2$ fixed points denoted by squares, and two $\mathbb{Z}_3$ fixed points, denoted by circles. Note that the three $\mathbb{Z}_2$ fixed points are interchanged by the $\mathbb{Z}_3$ action, while the two $\mathbb{Z}_3$ fixed points are images under the $\mathbb{Z}_2$ action.
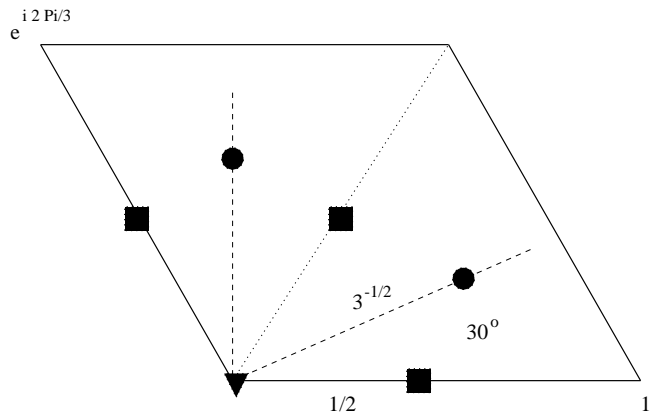

Figure 1. The $\mathbb{Z}_6$ (triangle), $\mathbb{Z}_3$ (circle), and $\mathbb{Z}_2$ (square) fixed points on one of the transverse two-tori. The "three-generation" $SU(5)$ model arises as the world volume theory of the 11 D3 branes, placed at the origin (the $\mathbb{Z}_6$ fixed point). The remaining 21 D3 branes must be distributed among the other fixed points or be at general positions, in a manner obeying some consistency conditions (see text).

Upon placing all 32 D3 branes at the origin, we can easily recover [8] the $SO(8) \times U(12)$ world volume theory of [10]. As these authors pointed out, it is not possible to continuously break the $SO(8) \times U(12)$ theory to an $SU(5)$ theory—in the D3 brane picture, as we will see below, this corresponds to the fact that it is impossible to smoothly move 21 of the 32 branes from the origin to the other fixed points; this is because the branes have to move in a $\mathbb{Z}_6$ symmetric manner,

i.e. in groups of 6. We now note that the twisted RR charge cancellation condition can be satisfied with a smaller number, n_0, of D3 branes at the origin; the general solution is: $n_0 = 8 + 3p, p = 0, 1, \ldots 8$. Since there are two distinct values of n_0 mod 6, it appears that there can be two disconnected branches of the moduli space. The minimal configuration with n_0=8 is smoothly connected to the n_0=32 configuration. The world volume theory has gauge group $U(4)$ and matter content consisting of three antisymmetric tensor representations (**6**'s) of $U(4)$, with no tree-level superpotential. This can be obtained by a continuous breaking of the $SO(8) \times U(12)$ world volume theory [10].

We now ask whether it is possible to have branches of the moduli space where an odd number of D3 branes is placed at the $\mathbb{Z}_2$ fixed points (i.e. the branch with $n_0 = 11$ mod 6, discussed above). Such brane configurations would correspond to a disconnected branch of moduli space—as discussed above, an odd number of D3 branes can not be continuously removed from the origin. There are, however, some nonperturbative consistency conditions that such configurations have to obey [12], [9]. These are best elucidated in terms of turning on Wilson lines in the T-dual nine-brane theory; for a more detailed discussion, see [8]. The nonperturbative consistency conditions have to do with the possibility to define spinors (which arise as nonperturbative states in the type-I theory) on the corresponding gauge bundles [12]. The conditions require that the number of D3 branes in the $\mathbb{Z}_2$ fixed points of each of the three two tori be even. This, however, still allows having an odd total number of branes removed from the origin. The simplest example involves only two of the three two-tori and allows us to put 9 D3 branes in the $\mathbb{Z}_2$ fixed points away from the origin is the following. We can place three (recall that they have to be placed in a $\mathbb{Z}_3$ symmetric manner) of the 9 D3 branes in the $\mathbb{Z}_2$ fixed point in both the first and second torus. This configuration clearly does not obey the above consistency condition, since the number of D3 branes at the $\mathbb{Z}_2$ fixed points in each two-torus is odd. We can satisfy the condition by placing three additional branes at the $\mathbb{Z}_2$

fixed points of the first torus only, while keeping them at the origin of the second, and three additional branes at the $\mathbb{Z}_2$ fixed points of the second torus, keeping them at the origin of the first. This configuration has a total of 9 branes at $\mathbb{Z}_2$ fixed points and obeys the consistency condition. It is easy to see, that similarly, by displacing branes in all the three two-tori, we can remove the desired 21 branes from the origin.

We can thus construct the interesting "three-generation" $SU(5)$ model: it arises upon placing 11 D3 branes at the origin. Performing the orientifold projection it is easy to see that the resulting $N = 1$ theory has the following matter content:

	$SU(5)$	$U(1)$
$A_{i=1,2,3}$	⊟	2
$\bar{Q}_{i=1,2,3}$	☐	-1

$$\text{(12)}$$

The theory has a renormalizable tree-level superpotential given by $W_{tree} = \epsilon^{ijk} A_i \bar{Q}_j \bar{Q}_k$.

We must distribute the remaining 21 D3 branes among the $\mathbb{Z}_2$ and $\mathbb{Z}_3$ fixed points, and general points of the torus (accounting for the various consistency conditions that need to be obeyed). For the branes at the $\mathbb{Z}_3$ fixed points, the world volume theories are obtained by only imposing the $\mathbb{Z}_3$ projection, while the theories on the branes at the $\mathbb{Z}_2$ fixed points are obtained by imposing only the orientifold projection. For branes at general points neither projection is required. The world volume gauge theories from these three sets of branes have different amounts of unbroken supersymmetries, since different projections are imposed. The $SO(2k_2 + 1)$ gauge theory from branes at the $\mathbb{Z}_2$ fixed points has $N = 4$ supersymmetry, as do the $U(1)$ gauge theories from branes at general points. The $[U(k_3)]^3$ gauge theory from branes at the $\mathbb{Z}_3$ fixed points has $N = 1$ supersymmetry.

So far we have ignored the effects of open strings stretched between branes at different fixed points. The lightest excitations of such strings are massive states which transform as fundamental-antifundamental under the respective world volume gauge groups. In the example we are considering, there can be two world volume theories with $N = 1$ supersymmetry. If supersymmetry were dynamically broken in one of these theories,

supersymmetry breaking would be communicated to the other gauge theory via the massive chiral multiplets just described (and, of course, by supergravity). A more precise investigation of this would probably involve considering details of the supersymmetry breaking dynamics and the stabilization of the dilaton [16]; we leave this for future work.

4. Supersymmetry breaking.

In this section, we will consider briefly the infrared dynamics of the $SU(5)$ model on the 11 D3 branes at the origin. We will show that, for fixed value of the string coupling and for any finite value of the orbifold blow-up parameter (i.e., the Fayet-Iliopoulos term of the anomalous $U(1)$), the ground state of the world-volume theory dynamically breaks supersymmetry. The analysis, details of which can be found in [8], relies on the fact that the dynamics of the $SU(5)$ three-generation model is known [14]. The theory (12) is an s-confining theory, and the low energy degrees of freedom are various composite meson and baryons:

$$
\begin{aligned}
C &= A \cdot \bar{Q} \cdot \bar{Q} & &\sim (\mathbf{3},\ \bar{\mathbf{3}})\ , \\
B &= A^5 & &\sim (\mathbf{6},\ \mathbf{1})\ , \\
M &= A^3 \cdot \bar{Q} & &\sim (\mathbf{8},\ \mathbf{3})\ ,
\end{aligned}
\tag{13}
$$

where we have shown their transformation properties under the global $SU(3)_A \times SU(3)_{\bar{Q}}$ symmetry. The confining superpotential is [14]:

$$
\begin{aligned}
W &= \frac{C_a^\alpha B^{\beta\gamma} M_\gamma^{\delta a} \epsilon_{\alpha\beta\delta} + M_\beta^{\alpha a} M_\gamma^{\beta b} M_\alpha^{\gamma c} \epsilon_{abc}}{\Lambda^9} \\
&+ \lambda \delta_\alpha^a C_a^\alpha,
\end{aligned}
\tag{14}
$$

where $a, b, ... (\alpha, \beta, ...)$ denote indices under the $SU(3)_{\bar{Q}(A)}$ symmetry, respectively, and the last term is the tree-level superpotential. The tree-level superpotential breaks the global symmetry to the diagonal $SU(3)_{diag}$ and lifts some of the classical flat directions (the moduli B and some of the M (13) are not lifted). The superpotential coupling λ in (14) is proportional to the value of the gauge coupling at the string scale (since the tree-level superpotential is the projection of the $N = 4$ superpotential).

One can now show [8] that (ignoring first the anomalous $U(1)$) the theory (12) with superpotential (14) has only a runaway supersymmetric solution (where the baryon $B \to \infty$). This suffices to show that the theory with the anomalous $U(1)$ breaks supersymmetry. To see this, note that the charges of all mesons and baryons under the $U(1)$ are positive. The D-term of the anomalous $U(1)$, therefore, generates a potential for the mesons and baryons that increases as their expectation value increases. Depending on the sign and value of the Fayet-Iliopolous term, this D-term potential could have a zero at finite values of the fields. However, the F-term potential, as we showed above, vanishes only at infinity. Therefore, one expects that supersymmetry is broken for generic values of the Fayet-Iliopoulos term; it is only possible to have vanishing D- and F-term potentials for infinite (negative) value of the FI term. For finite values of the orbifold blow-up parameter, however, supersymmetry is always broken.

We also note that if we consider the branch of moduli space with 8 mod 6 D3 branes at the origin, i.e. the $SU(4)$ theory, we would find that the theory has restored supersymmetry for finite values of the blow-up parameter (such that the one-loop Fayet-Iliopoulos term is cancelled). This is because, in contrast to the $SU(5)$ model, the $SU(4)$ theory has a branch of moduli space where no dynamical superpotential is generated—this can be inferred from [15] by noting that the $SU(4)$ theory with three $\mathbf{6}$'s is equivalent to the $SO(6)$ theory with three vectors. The breaking of supersymmetry is then purely D-term and vanishes for vanishing FI parameter.

We note that in the above discussion of supersymmetry breaking we ignored the closed string modes. Once these are allowed to fluctuate, there are, as usual, runaway directions along which supersymmetry is restored. We leave a detailed investigation of this for future work.

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